THEORETICAL ISSUES IN NEUTRINO PHYSICS

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I review a number of the open questions about neutrino properties, critique recent hints of neutrino mass, and discuss one recently proposed neutrino mass matrix to illustrate the direction in which we may be headed. I also present one example of the implications of these new developments for astrophysics.

1   Introduction

In this talk I would like to discuss a number of the open questions we have entertained about neutrinos since they were first postulated more than 65 years ago, as well as a few of the possible answers that may result from atmospheric, solar, and terrestrial neutrino experiments. Indeed, the list of open questions has proven surprisingly resistant to experiment:

• What are the masses of the ν_e, ν_μ, and ν_τ, and what will these masses tell us about new scales beyond the standard model?

• Are there massive neutrinos beyond the range tested by measurements of the Z^0 width?

• What are the particle-antiparticle conjugation properties of the known neutrinos?

• Do neutrinos have nonzero electromagnetic moments (magnetic dipole, electric dipole, or anapole), or perhaps a nonzero charge radius?

• Do neutrinos of different flavor mix to produce neutrino oscillations?

• Can we prove cosmic background neutrinos exist? What role have they played in determining the present structure of our universe? Do they comprise an appreciable fraction of the dark matter?

• What is the role of neutrinos in core-collapse supernovae? Are nonstandard neutrino properties essential to the explosion mechanism or to the associated nucleosynthesis (r-process, ν-process)?

• What are the sources of very high energy neutrinos in astrophysics? Are they associated with gamma ray bursts or active galactic nuclei? Does the standard model properly describe their propagation through and interactions with the cosmic background radiation?

The above list could be extended for several more pages. What we do know about neutrinos is, in some respects, equally puzzling. For example direct

mass limits tell us that

$$\text{m}(\nu_e) \lesssim (3-5) \text{ eV } [1]$$
$$\text{m}(\nu_\mu) \lesssim 170 \text{ keV } [2]$$
$$\text{m}(\nu_\tau) \lesssim 18.2 \text{ MeV } [2]$$

Naively, such small values pose a problem for theorists hoping to extend the standard model by unifying the known particles into larger multiplets. For example, an attractive idea [3] might be multiplets containing all of the particles of a given family

$$\begin{pmatrix} \nu \\ e \end{pmatrix} \xrightarrow[\substack{grander \\ model}]{e.g.} \begin{pmatrix} u \\ d \\ \nu \\ e \end{pmatrix}$$

One then might expect (to within group theory factors) that the members of the multiplet would couple to the mass-generating fields in a similar way, and thus have about the same mass, $m_i \sim o(m_D)$. Now the u and d quarks and the electron have masses on the order of an MeV, but the ν_e clearly breaks the pattern: it is at least six orders of magnitude lighter. One popular resolution of this dilemma is connected with additive quantum numbers. For example
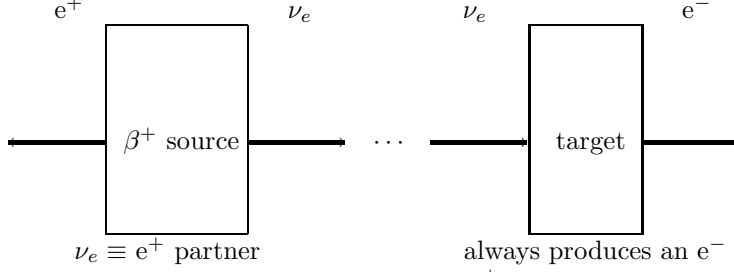
$$e^- \rightarrow e^+$$

under charge conjugation, clearly producing an orthogonal antiparticle, distinguished from the electron by an additive quantum number (the charge). However the corresponding question for the ν — does there exist a distinct antiparticle, $\bar{\nu}$? — is not so easy to answer.

Before tackling this question, it is helpful to discuss some of the complicating issues connected with the handedness of massive neutrinos. Consider a massive left-handed neutrino moving at a velocity $v < c$. Now boost the observer to frame moving faster than that neutrino's velocity
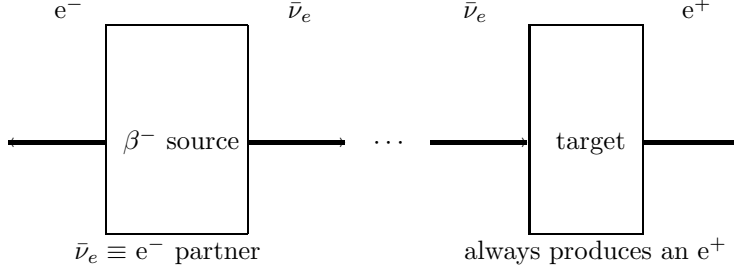
$$\nu_{LH} \begin{matrix} \overleftarrow{s} \\ \overrightarrow{p} \end{matrix} \xrightarrow[boost]{} \begin{matrix} \overleftarrow{s} \\ \overleftarrow{p} \end{matrix} \nu_{RH}$$

In the boosted frame the neutrino is right-handed. From this exercise we learn that the Lorentz structure of any model describing massive neutrinos demands both ν_{LH} and ν_{RH}. As the interactions of the standard model are V-A, a closely associated point is that neutrino masses break the "γ_5" invariance of interactions, leading to interesting effects of order (m_ν/E_ν): we will see an illustration of this in a later discussion of double beta decay.

It would seem that the logical way to resolve the issue of a $\bar{\nu}$ distinguishable from the ν is to test the properties of these particles experimentally. If neutrinos are produced by a β^+ source and their interactions tested in a target downstream, one finds

$$\text{e}^+ \qquad\qquad \nu_e \qquad\qquad \nu_e \qquad\qquad \text{e}^-$$

$$\nu_e \equiv \text{e}^+ \text{ partner} \qquad\qquad \text{always produces an e}^-$$

That is, if we define a ν_e as the partner of the e$^+$ in a β decay reaction, then we observe that ν_es always produce e^-s when they react in a target. Similarly

$$\text{e}^- \qquad\qquad \bar{\nu}_e \qquad\qquad \bar{\nu}_e \qquad\qquad \text{e}^+$$

$$\bar{\nu}_e \equiv \text{e}^- \text{ partner} \qquad\qquad \text{always produces an e}^+$$

where the $\bar{\nu}_e$ has been defined as the partner of the e^- in β decay. It would appear that the ν_e and $\bar{\nu}_e$ so defined are operationally distinct

$$\nu_e\text{s produce e}^-\text{s}$$
$$\bar{\nu}_e\text{s produce e}^+\text{s}$$

This motivates the introduction of a distinguishing quantum number (lepton number)

$lepton$	l_e
e$^-$	$+1$
e$^+$	-1
ν_e	$+1$
$\bar{\nu}_e$	-1

If we require that lepton number is additively conserved,

$$\sum_{in} l_e = \sum_{out} l_e,$$

the "experimental" results discussed above then follow.

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The experiments described above are done by nature virtually in the process of neutrinoless $\beta\beta$ decay, $(A,Z) \rightarrow (A,Z+2) + e^- + e^-$, as depicted in Fig. 1. The illustrated process cannot take place if the emitted neutrino is orthogonal to the antineutrino that must be absorbed on the second nucleon

$$\nu_e(l = +1) \perp \bar{\nu}_e(l = -1).$$

Results for the current generation of enriched ^{76}Ge $\beta\beta$ decay searches [4]

$$\tau_{1/2}(\mathrm{Ge}^{76}) \gtrsim 2 \cdot 10^{25}\mathrm{y}$$

imply a limit on the electron neutrino Majorana mass of $\langle \mathrm{m}_\nu^{\mathrm{Maj}} \rangle \lesssim 0.4$ eV, where [5]

$$\langle \mathrm{m}_\nu^{\mathrm{Maj}} \rangle = \sum_{i=1}^{2n} \eta_i^{\mathrm{CP}} \mathrm{U}_{ei}^2 \mathrm{m}_i.$$

Here m_i is the mass of the ith eigenstate, U_{ei} is the amplitude of that mass eigenstate in $|\nu_e\rangle$, and η_i^{CP} is the relative CP of the ith mass eigenstate. Thus CP conservation has been assumed.
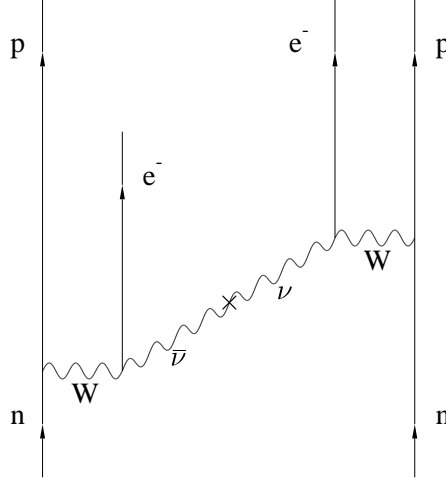

Figure 1. Two-nucleon diagram for neutrinoless $\beta\beta$ decay. The amplitude vanishes if the ν and $\bar{\nu}$ are distinct, i.e., carry different lepton numbers.

While the nonobservation of neutrinoless $\beta\beta$ decay is consistent with a Dirac neutrino—one where the ν and $\bar{\nu}$ are distinguished by their opposite lepton numbers—it is not required, due to the V-A character of standard model weak interactions. The replacements

$$\nu_e \rightarrow \nu_e^{LH}$$

$$\bar{\nu}_e \rightarrow \nu_e^{RH}$$

lead to a helicity missmatch in the above diagram and a decay rate suppressed by $o(\frac{\langle m_k^{Maj} \rangle}{E_\nu})^2$. Thus the absence of neutrinoless $\beta\beta$ decay is consistent with a Majorana neutrino ($\nu = \bar{\nu}$) if the electron neutrino Majorana mass is light, as indicated above.

2   The Neutrino Mass Matrix [5,6]

A Majorana ν_e corresponds to the limiting case where a state of definite mass has two components, with both the boosts and particle-antiparticle conjugation (CPT, or CP in the limit of CP conservation) coupling one component to the other, as depicted below. A Dirac ν_e is a four-component neutrino, where the ν and $\bar{\nu}$ are distinguished by their lepton numbers, and thus where the boosts and CP/CPT operations connect distinct components.
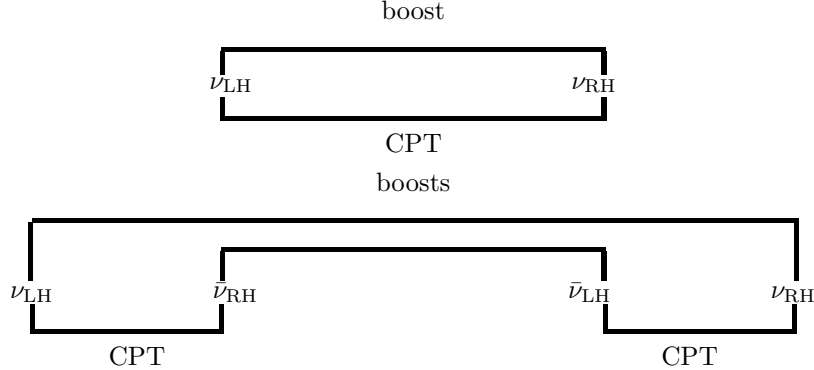
We now proceed through a simple exercise of generalizing these limiting cases to one where several mass eigenstates may contribute, and where both Dirac and Majorana mass terms arise. The starting point is a Dirac field from which we project the four components, using the R/L and charge conjugation projection operators:

$$\psi_{R/L} = \tfrac{1}{2}(1 \pm \gamma_5)\psi$$
$$\mathrm{C}\psi_{R/L}\mathrm{C}^{-1} = \psi_{R/L}^c$$

We apply this to the mass term for the Dirac equation

$$\mathcal{L}_m(x) \sim m_D \bar{\psi}(x)\psi(x) \Rightarrow M_D \bar{\Psi}(x)\Psi(x)$$

where m_D has been replaced by a nondiagonal 3×3 matrix M_D in flavor space and

$$\Psi = \begin{pmatrix} \psi^e \\ \psi^\mu \\ \psi^\tau \end{pmatrix}$$

The resulting mass matrix

$$(\bar{\Psi}_L^c, \bar{\Psi}_R, \bar{\Psi}_L, \bar{\Psi}_R^c) \begin{pmatrix} 0 & 0 & 0 & M_D^T \\ 0 & 0 & M_D & 0 \\ 0 & M_D^\dagger & 0 & 0 \\ M_D^* & 0 & 0 & 0 \end{pmatrix} \begin{pmatrix} \Psi_L^c \\ \Psi_R \\ \Psi_L \\ \Psi_R^c \end{pmatrix}$$

then allows for flavor oscillations, as M_D is assumed to be nondiagonal.

While the upper left and lower right quadrants of this matrix must be zero because the left- and right-handed projectors annihilate each other, obviously additional terms can be introduced elsewhere if we respect the requirement of hermiticity. Specifically,

$$\mathcal{L}_m(x) \Rightarrow M_D \bar{\Psi}(x)\Psi(x) + \bar{\Psi}_L^c M_L \Psi_L + \bar{\Psi}_R^c M_R \Psi_R$$

so that the mass matrix becomes

$$(\bar{\Psi}_L^c, \bar{\Psi}^R, \bar{\Psi}_L, \bar{\Psi}_R^c) \begin{pmatrix} 0 & 0 & M_L & M_D^T \\ 0 & 0 & M_D & M_R^\dagger \\ M_L^\dagger & M_D^\dagger & 0 & 0 \\ M_D^* & M_R & 0 & 0 \end{pmatrix} \begin{pmatrix} \Psi_L^c \\ \Psi_R \\ \Psi_L \\ \Psi_R^c \end{pmatrix}$$

The new Majorana mass terms break the local gauge invariance $\psi(x) \rightarrow e^{i\alpha(x)}\psi(x)$ associated with a conserved lepton number. It is these nonDirac mass terms that can generate the nonzero $\langle m_\nu^{Maj} \rangle$ that gives rise to neutrinoless $\beta\beta$ decay.

One can proceed to diagonalize this matrix

$$\Psi_{\nu(e)}^L = \sum_{i=1}^{2n} U_{ei}^L \tilde{\nu}_i(x) \ \text{ with masses } m_i$$

The eigenstates are two-component Majorana neutrinos [5], yielding the proper $2 \times 2n = 4n$ degrees of freedom, where n is the number of flavors. We can recover the Majorana and Dirac limits:

• If $M_R = M_L = 0$, the eigenstates of this matrix become pairwise degenerate, allowing the $2n$ two-component eigenstates to be paired to form n four-component Dirac eigenstates.

• If $M_D = 0$, the left- and right-handed components decouple, yielding n left-handed Majorana eigenstates with standard model interactions.

There are interesting physical effects associated with these limits. Dirac neutrinos can have magnetic dipole, electric dipole (CP and T violating), and anapole (P violating) moments, as well as nonzero charge radii. Majorana neutrinos can have anapole moments but only transition magnetic and electric dipole moments. Yet transition moments are quite interesting in the context of matter-enhanced spin-flavor oscillations [7]. The most stringent limits on both diagonal and transition magnetic and electric dipole moments come from red giant evolution, where the enhanced neutrino pair production delays core He ignition. This yields [8]

$$|\mu_{ij}| \lesssim \text{few} \cdot 10^{-12} \mu_B$$

a bound that is approximately two orders of magnitude more restrictive than the laboratory limit.

Neutrinos are unique among the fermions in allowing both Dirac and Majorana mass terms, a consequence of the absence of any obvious additive quantum numbers that must change under particle-antiparticle conjugation. The presence of both mass terms provides an attractive explanation for small neutrino masses, the seesaw mechanism of Gell-Mann, Ramond, Slansky, and Yanagida [9]. As $\beta\beta$ decay suggests a left-handed Majorana mass much smaller than typical Dirac masses, while right-handed interactions are not seen at low energies and thus might be characterized by mass scales well beyond the standard model, the following mass matrix is natural:

$$\begin{pmatrix} 0 & m_D \\ m_D & m_R \end{pmatrix} \Rightarrow m_\nu^{light} = m_D(\frac{m_D}{m_R})$$

Thus m_D/m_R is the needed small parameter explaining why neutrinos are so much lighter than their charged partners. If the ν_τ mass is on the order of 0.1 eV (a value suggested by atmospheric neutrinos), and the Dirac mass is taken from $m_{top} \sim 180$ GeV, this yields $m_R \sim 0.3 \times 10^{15}$ GeV, a value reasonably close to the GUT scale, $M_{GUT} \sim 10^{16}$. The massive right-handed neutrino of the seesaw fits naturally into various extended models [3,10]. For example, the

16-dimensional family multiplet of SO(10) is

$$(\vec{u}_L, \vec{d}_L, \vec{u}_R, e_R, \vec{d}_R, e_L, \nu_L, \nu_R).$$

That is, the assignment provides a natural spot to be filled by a heavy, chargeless, right-handed neutrino. Thus the key question is whether current hints of neutrino mass are telling us about physics at 10^{15} GeV.

3 Handicapping the Hints of Mass [11]

The solar neutrino problem, the discrepancy between the predictions of the standard solar model (SSM) [12] and the results of the ^{37}Cl, GALLEX and SAGE, and KamiokaII/III and SuperKamiokande experiments, was described by Hamish Robertson. There are two somewhat distinct arguments that this discrepancy requires new particle physics:
• There is the $\sim 3\sigma$ argument based on global fits to the various experiments, using undistorted neutrino spectra but making no other assumptions, or only weak assumptions (such as a steady state model where the luminosity constrains present fusion rates), about the solar model (see, for example, [13]). In such fits, an unphysical result is obtained, a negative ^{7}Be flux.
• There is also a $\sim 5\sigma$ argument in which the general temperature dependence of standard and nonstandard models is used. This requires some additional assumptions, such as solar burning under the conditions of chemical equilibrium [14], but is still independent of many details of the solar model. The conclusion from experiment that $\phi(^8\text{B}) \sim 0.4\phi^{SSM}(^8\text{B})$ and from solar models that $\phi(^8\text{B}) \propto T_c^{18}$, where T_c is the core temperature, leads to the conclusion

$$T_c \sim 0.96 T_c^{\text{SSM}}$$

That is, a cool solar core is required. Consequently as

$$\Phi(^7\text{Be})/\Phi(^8\text{B}) \sim T_c^{-10}$$

such a cooler core requires

$$\frac{\Phi(^7\text{Be})}{\Phi(^8\text{B})} \sim 1.5 \frac{\Phi^{\text{SSM}}(^7\text{Be})}{\Phi^{\text{SSM}}(^8\text{B})}$$

But experimentally we find $\Phi(^7\text{Be}) \sim 0$, so that the experimental ratio is low, as would be expected for a hotter core. This is the crux of the second argument: experiment leads to a contradiction in the absence of new particle physics, with one observation ($\Phi(^8\text{B})$) requiring a cooler core and a second ($\Phi(^7\text{Be})/\Phi(^8\text{B})$) a hotter one. This conflict is nicely illustrated in Fig. 2,
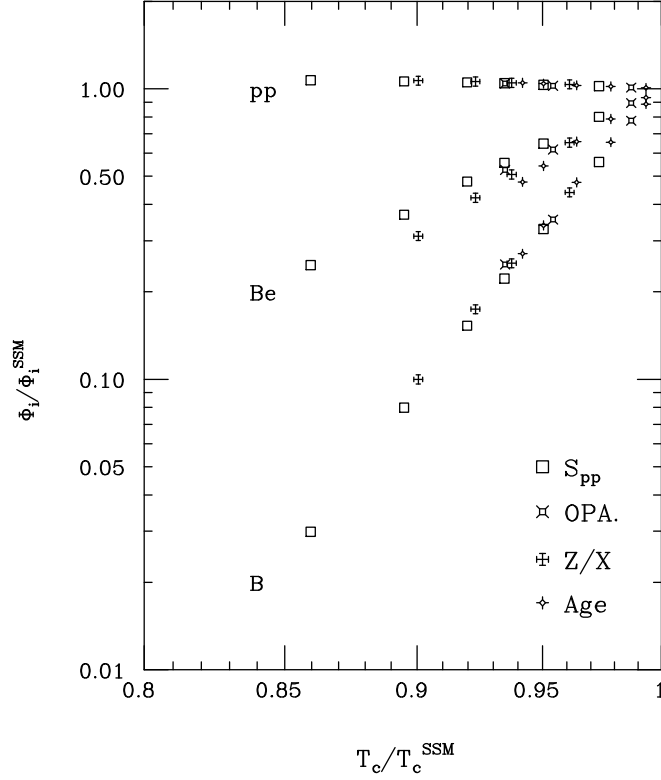
taken from Castellani et al. [15].

Figure 2. The response of the pp, ^{7}Be, and ^{8}B solar neutrino fluxes to the indicated variations in solar model input parameters, displayed as a function of the resulting central temperature T_c. From [15].

These two indirect arguments are quite convincing, leading many in the field to conclude that the solar neutrino problem demands new particle physics. Yet, as the conclusions are based on combining the results from several experiments, no one of which by itself requires new physics, it would be very reassuring to see a "smoking gun" signal. The current generation of active detectors - SuperKamiokande with its spectral sensitivity and SNO with its direct sensitivity to neutral currents - has the potential to yield such a signal. Currently the SuperK spectral distribution does show some structure, but

of an unexpected kind: there is a surprising excess of high-energy electrons, similar to those expected from the ^{3}He + p reaction if the standard cross section estimate were to be multiplied by about a factor of 30 [16].

Experiments sensitive to atmospheric neutrinos have traditionally expressed their results in terms of a ratio of ratios

$$R = \frac{(N_\mu/N_e)_{DATA}}{(N_\mu/N_e)_{MC}}$$

determined from the measured and calculated (with Monte Carlo codes) muon-like and electron-like neutrino rates. Robertson's talk gave many of the reasons the atmospheric neutrino problem is considered such a convincing argument for new physics:
•The SuperK ratio has been very accurately determined, $\sim 0.61 \pm 0.03 \pm 0.05$.
•There is good consistency between sub-GeV/multi-GeV and fully contained/partially contained data sets.
• The change in the ratio with zenith angle provides direct evidence for neutrino oscillations.
• The results for R are consistent among the four detectors with the largest data sets (SuperK, SoudanII, IMB, Kamiokande).
• The up-down difference is "self-normalizing," almost independent of the calculated atmospheric fluxes.

The favored interpretation of the SuperK and other atmospheric neutrino results is a large-mixing-angle $\nu_\mu \rightarrow \nu_\tau$ oscillation. There have been some questions raised about this interpretation:
• The absolute rates can be fit as well by an excess of e-like events as by a deficit in μ-like events.

	data	MonteCarlo
$e-like$	983	812
$\mu-like$	900	1218

•The SuperK $\sin^2 2\theta - \delta m^2$ region favors smaller values of δm^2 than those found by Kamioka and SoudanII, though there is a region of overlap.
• There is some tension between the SuperK shape fit and R, with the results for R largely agreeing with other atmospheric neutrino experiments, but lying mostly outside the region favored by the shape fit.
• The results require very large mixing angles, with more than half of the 90% confidence level region lying in the unphysical region ($\sin^2 2\theta$ exceeding 1) in an unconstrained fit.

Yet despite these concerns, the most striking aspect of the SuperK results is the azimuthal dependence, which is direct evidence for neutrino oscillations. And, as will be illustrated below, a scenario with maximal mixing angles can be made to fit nicely with other hints of neutrino mass.

There is a third hint of neutrino mass, direct laboratory evidence from the LSND experiment, although the allowed region is mostly excluded by the KARMEN experiment. It is very difficult to evaluate this situation. If LSND were proven to be correct, its inclusion with the solar and atmospheric neutrino problems very much strains three-flavor fits to these three results.

4   Prejudices and One Possible Pattern

To provide some picture of how these various results might fit together to form some pattern, I now discuss a recent paper by Georgi and Glashow [17]. The assumptions of their construction are:
• Three light Majorana neutrinos
• The atmospheric neutrino problem is due to $\nu_\mu \rightarrow \nu_\tau$ oscillations, since the $\nu_\mu \rightarrow \nu_e$ alternatively is ruled out by the Chooz experiment.
•This oscillation is nearly maximal with $\sin 2\theta_{23} \sim 1$ and $5 \cdot 10^{-4}$ eV$^2 \lesssim \delta m_{23} \lesssim 6 \cdot 10^{-3}$ eV^2.
•The solar neutrino problem is due to oscillations with $6 \cdot 10^{-11}$ eV$^2 \lesssim \delta m^2 \lesssim 2 \cdot 10^{-5}$ eV^2.
• The neutrino masses are constrained to satisfy $m_1 + m_2 + m_3 \sim 6$ eV in order to generate hot dark matter for large scale structure formation (a somewhat speculative condition).
•The absence of neutrinoless double β decay requires $\langle m_\nu^{Maj} \rangle \lesssim 0.4$eV, so choose $\langle m_\nu^{Maj} \rangle \sim 0$.
•Because of the LSND/KARMEN conflict, the LSND results are not considered.

These constraints lead to a pattern of three nearly degenerate massive neutrinos with $m_i \sim M$ and a simple mass matrix that accounts for the atmospheric and solar neutrino problems through vacuum oscillations,

$$\mathcal{M}^L = M \begin{pmatrix} 0 & \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} \\ \frac{1}{\sqrt{2}} & \frac{1}{2} & \frac{-1}{2} \\ \frac{1}{\sqrt{2}} & \frac{-1}{2} & \frac{1}{2} \end{pmatrix} \begin{matrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{matrix}$$

That is,

$$|\nu_e\rangle = \overbrace{\frac{1}{\sqrt{2}}|\nu_1\rangle - \frac{1}{\sqrt{2}}|\nu_2\rangle}^{\Delta \sim \delta m^2_{solar}/2M} \ (opp.CP)$$

$$|\nu_\mu\rangle = \frac{1}{2}|\nu_1\rangle + \frac{1}{2}|\nu_2\rangle + \frac{1}{\sqrt{2}}|\nu_3\rangle$$

$$|\nu_\tau\rangle = \frac{1}{2}|\nu_1\rangle + \underbrace{\frac{1}{2}|\nu_2\rangle - \frac{1}{\sqrt{2}}|\nu_3\rangle}_{\Delta \sim \delta m^2_{atmos}/2M}$$

where the two mass eigenstates comprising the ν_e interfere in the $\beta\beta$ decay mass because they have opposite CP.

This scenario has the following consequences:
• The $\beta\beta$ decay mass and solar neutrino mass δm^2 are related by $\langle m_\nu^{Maj}\rangle = \delta m^2_{solar}/4M$, with the resulting rates for $\beta\beta$ decay thus being very small.
• The $\nu_\mu \to \nu_\tau$ oscillation is maximal over terrestrial distances as ν_3 beats against ν_1 and ν_2; the small splitting of the latter two mass eigenstates leads to $|\nu_e\rangle \to (|\nu_\mu\rangle + |\nu_\tau\rangle)$ oscillations only over larger distances. This solar neutrino oscillation is also maximal.
• Interestingly, the MSW mechanism is not used.
This kind of mass matrix can arise naturally in model schemes, as has been shown recently by Mohapatra and Nussinov [18]. Clearly it is just one possibility among many, but suggests that the hints of massive neutrinos we now have may yet conform to a simple pattern.

5   Neutrinos and the r-process

As this talk is at an end, let me mention very quickly some interesting connections between massive neutrinos and the explosive stellar synthesis of heavy nuclei by the rapid-neutron-capture or r-process. About half of the heavy elements above the iron group are thought to be created by this process, where neutron capture is faster than β decay, so that the usual weak equilibrium condition of nuclei is replaced by $(n, \gamma) \leftrightarrow (\gamma, n)$ equilibrium. Thus the synthesis occurs along a path through very neutron rich nuclei near

the neutron drip line.

This process requires extraordinarily explosive conditions
$$\rho(n) \sim 10^{20} \text{ cm}^{-3} \quad T \sim 10^9 K \quad t \sim 1 \text{sec}$$
Probably the most plausible of the conjectured sites for the r-process is the high-entropy, neutron rich gas near the mass cut of a Type II supernova, the last material to be ejected. As this material expands off the proto-neutron star, it undergoes an alpha-rich freezeout, and then an alpha process that may continue to nuclei near $A \sim 100$. The result is a soup of αs, a few heavy seed nuclei, and excess neutrons. While detailed modeling of this "hot bubble r-process" fails in some details, the basic requirement of ~ 100 neutrons per seed nucleus appears achievable [19].

This matter experiences an enormous fluence of neutrinos, emitted by the cooling protoneutron star. As weak equilibrium in maintained among the various neutrino species through most of their random walk out of the neutron star, there is an approximate equipartion of energy per flavor. However, the location of the neutrinosphere (the surface of last scattering) does depend on flavor because of the strong, charged current interactions of the ν_es and $\bar{\nu}_e$s. Neutrinos that decouple earlier do so at the higher ambient temperatures characterizing the smaller neutrinospheres. The results are
$$T(\nu_\mu, \nu_\tau) \sim 8 \text{ MeV}$$
$$T(\bar{\nu}_e) \sim 4.5 \text{ MeV}$$
$$T(\nu_e) \sim 3.5 \text{ MeV}$$
That is, the heavy flavor neutrinos are expected to be, on average, significantly more energetic than the ν_es and $\bar{\nu}_e$s.

For the usual seesaw pattern of neutrino masses and a cosmological interesting ν_τ (i.e., a heavy neutrino with a mass in the neighborhood of 10 eV), the full MSW pattern is shown in Fig. 3. If the $\nu_e - \nu_\mu$ crossing is responsible for the solar neutrino problem, a second crossing, $\nu_e - \nu_\tau$, is expected at a density large compared to that of the solar core, but small compared to the location of the supernova neutrinosphere ($\sim 10^{12}$ g/cm^3). For a very large range of mixing angles, this crossing is adiabatic and thus leads to $\nu_e \leftrightarrow \nu_\tau$ conversion. These spectra thus change identities, leading to an anomalously hot ν_e flux from a Type II supernova.

As the ν-nucleon cross section is proportional to E_ν^2,
$$\nu_e + n \rightarrow e^- + p \quad \text{is enhanced}$$
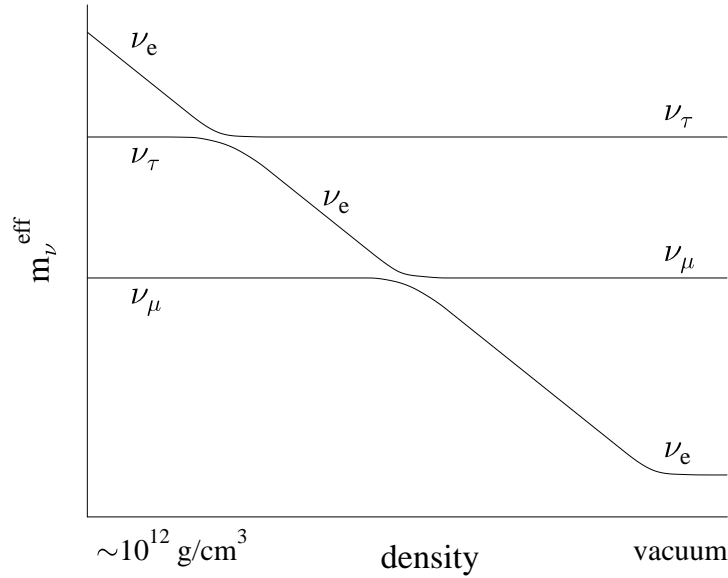$$\bar{\nu}_e + p \rightarrow e^+ + n \quad \text{is unchanged}$$

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Figure 3. Three-flavor MSW diagram showing possible ν_e–ν_μ solar and ν_e–ν_τ supernova crossings.

For a rather extensive range of $\nu_e \leftrightarrow \nu_\tau$ mixing angles and δm^2, this crossing then destroys the r-process: the hotter ν_es drive the matter proton rich [20]. Thus, if one accepts this location as the site of the r-process, very strong constraints on cosmologically interesting ν_τs are obtained.

This provides a nice closing for this talk. I began by discussing how neutrino masses and other neutrino properties might tell us about physics beyond the standard model - including, perhaps, physics at the GUT scale. This ending shows how the pattern of neutrino masses may be equally relevant to our low-energy world: the formation of large scale structure and the existence of the transuranic elements may be issues connected by the masses and mixing angles of neutrinos.

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